

Diphoton excess from hidden U(1) gauge symmetry with large kinetic mixing

Fuminobu Takahashi^{a,b}, Masaki Yamada^a and Norimi Yokozaki^a

^a *Department of Physics, Tohoku University,
Sendai, Miyagi 980-8578, Japan*

^b *Kavli IPMU (WPI), UTIAS,
The University of Tokyo, Kashiwa, Chiba 277-8583, Japan*

Abstract

We show that the 750 GeV diphoton excess can be explained by introducing vector-like quarks and hidden fermions charged under a hidden U(1) gauge symmetry, which has a relatively large coupling constant as well as a significant kinetic mixing with U(1)_Y. With the large kinetic mixing, the standard model gauge couplings unify around 10^{17} GeV, suggesting the grand unified theory without too rapid proton decay. Our scenario predicts events with a photon and missing transverse momentum, and its cross section is related to that for the diphoton excess through the kinetic mixing. We also discuss other possible collider signatures and cosmology, including various ways to evade constraints on exotic stable charged particles. In some cases where the 750 GeV diphoton excess is due to diaxion decays, our scenario also predicts triphoton and tetraphoton signals.

1 Introduction

The diphoton excess with an invariant mass around 750 GeV was recently reported by the ATLAS [1] and CMS [2] collaborations; for a spin-0 particle with a narrow width approximation, the local significance is estimated to be 3.9σ and 2.9σ , respectively. While more data is certainly needed to confirm if the signal is real or just a statistical fluke, its high statistical significance in the clean analysis using photons triggered enthusiasm and exuberance for the new physics beyond the standard model (SM), followed by the appearance of many theoretical papers.

Among various models proposed so far, the simplest one is to include a gauge singlet (pseudo)scalar coupled to vector-like quarks and/or leptons (see e.g. Refs. [3–5] for the early works).^{#1} In this model, the (pseudo)scalar is produced via gluon fusion and decays into a pair of photons through one-loop diagrams with the extra quarks/leptons running in the loop. The diphoton excess can be explained if the product of the production cross section times branching ratio to two photons is in the range of $5 - 10$ fb. This gives a preference to a relatively large branching fraction to diphotons, which necessitates either multiple extra matter fields and/or large hypercharges ($Y \gtrsim 1$) of the extra field running in the loop.

In this paper we consider a possibility that the large hypercharges are originated from unbroken hidden $U(1)_H$ gauge symmetry which has a relatively strong gauge coupling and a significant kinetic mixing with $U(1)_Y$. Then, hidden fermions acquire large hypercharges due to the kinetic mixing, and the induced hypercharges are generically irrational. We will show that the diphoton excess can be explained by the (pseudo)scalar coupled to gluons and photons through the extra quark/hidden fermion loop diagrams.

Our scenario is based on a rather simple $U(1)_H$ extension of the standard model, which enables us to make a definite prediction that can be tested soon at the LHC Run-2. Since the hidden fermions are charged under $U(1)_H$, the (pseudo)scalar responsible for the diphoton excess can also decay into $\gamma\gamma'$, where γ' denotes the hidden photon. Thus, our scenario predicts events with a photon and missing momentum,^{#2} and we will see that its production cross section times branching fraction is simply related to that for the diphoton excess through the kinetic mixing. The events with a photon and missing momentum have been searched for at the LHC Run-1 [11, 12] and Run-2 [13], and there is an upper bound on the production cross section. We will see that the experimental bound places a *lower* bound on the kinetic mixing.

A large kinetic mixing with $U(1)_H$ is known to modify the normalization of the hypercharge so that the gauge coupling unification is improved [14]. We will show that this is indeed the case in our model, taking account of contributions of the extra matter fields to the renormalization group (RG) equations.^{#3} The hidden fermions acquire hypercharges

^{#1} It is still a puzzle why such a (pseudo)scalar coupled to gluons and photons exists in nature. One possible answer is to relate it to the QCD axion (or its bosonic partner, saxion) which solves the strong CP problem [6–9].

^{#2} See Ref. [10] for a related work.

^{#3} In supersymmetric models with the grand unification, the diphoton excess may indicate the light gluino

through the kinetic mixing, and they are cosmologically stable. Such stable exotic charged particles, if produced abundantly in the early Universe, could affect the big bang nucleosynthesis (BBN) [17–19] as well as cosmic microwave background radiation (CMB) [20, 21]. Also there are various experimental searches for exotic fractional or multi charged particles [22–24]. We will discuss several possibilities to evade those constraint. Finally, we will discuss other possible “diphoton” excesses at different energies if the 750 GeV diphoton excess is due to diaxion decays.

The rest of this paper is organized as follows. In Sec. 2 we show that the diphoton excess can be explained by introducing a gauge singlet (pseudo)scalar and vector-like quarks and hidden fermions, the latter of which is charged under $U(1)_H$. In Sec. 3 we study the gauge coupling unification in the presence of the large kinetic mixing. Cosmological implications are discussed in Sec. 4. The last section is devoted for conclusions.

2 Kinetic mixing with hidden $U(1)$

2.1 Preliminaries

Let us first quickly review the effect of a kinetic mixing between two $U(1)$ s, $U(1)_1$ and $U(1)_2$. We will shortly apply the results to the kinetic mixing between $U(1)_Y$ and a hidden $U(1)$ gauge symmetry, $U(1)_H$.

Let us consider the Lagrangian [28],

$$\mathcal{L} = -\frac{1}{4}F_1^{\prime\mu\nu}F'_{1\mu\nu} - \frac{1}{4}F_2^{\prime\mu\nu}F'_{2\mu\nu} - \frac{\chi}{2}F_1^{\prime\mu\nu}F'_{2\mu\nu}, \quad (2.1)$$

where $F_i^{\prime\mu\nu} \equiv \partial^\mu A_i^{\prime\nu} - \partial^\nu A_i^{\prime\mu}$ ($i = 1, 2$) represent the field strength of $U(1)_i$, and χ is the kinetic mixing between them. The kinetic mixing can be removed by the following transformation,

$$A_1^{\mu'} = \frac{A_1^\mu}{\sqrt{1-\chi^2}}, \quad (2.2)$$

$$A_2^{\mu'} = A_2^\mu - \frac{\chi}{\sqrt{1-\chi^2}}A_1^\mu, \quad (2.3)$$

where A_1^μ and A_2^μ are canonically normalized gauge fields. Hereafter we call this new basis ($A_i^{\mu'}$) as the canonical basis to distinguish it from the original basis (A_i^μ). In the canonical basis, gauge couplings e_1 and e_2 are written in terms of the kinetic mixing χ and the gauge couplings in the original basis e'_1 and e'_2 such as

$$e_1 = \frac{e'_1}{\sqrt{1-\chi^2}}, \quad (2.4)$$

$$e_2 = e'_2. \quad (2.5)$$

of 2–3 TeV, which originates from changes of RG equations with extra matter fields [15].

In the canonical basis, any matter fields charged under $U(1)_1$ in the original basis are still coupled to A_1^μ with a rescaled gauge coupling, e_1 . On the other hand, the matter field charged under $U(1)_2$ acquires an induced charge of $U(1)_1$ in the canonical basis. For instance,

$$q_2 e_2' \bar{\psi} \gamma^\mu \psi A_{2\mu}' = q_2 e_2 \bar{\psi} \gamma^\mu \psi A_{2\mu} - \frac{\chi}{\sqrt{1-\chi^2}} q_2 e_2 \bar{\psi} \gamma^\mu \psi A_{1\mu}. \quad (2.6)$$

Thus, through the kinetic mixing, a matter field with a charge q_2 under $U(1)_2$ acquires a charge, $-\frac{\chi}{\sqrt{1-\chi^2}} \frac{q_2 e_2}{e_1}$, under $U(1)_1$ in the canonical basis. The induced charge is generically irrational, and can be larger than unity depending on the relative size of the gauge couplings and the kinetic mixing.

2.2 Diphoton excess

Now we apply the above result to $SU(2)_L \times U(1)_Y \times U(1)_H$, where we assume a kinetic mixing χ between $U(1)_Y$ and $U(1)_H$. Suppose that there is a hidden matter field ψ with a charge q_H under $U(1)_H$. Then, in the canonical basis, the hidden matter field acquires an electric charge,

$$q_{\text{eff}} = -\frac{\chi}{\sqrt{1-\chi^2}} \frac{q_H e_H}{e_{\text{EM}}} \cos \theta_w, \quad (2.7)$$

where e_{EM} and e_H are gauge couplings of $U(1)_{\text{EM}}$ and $U(1)_H$, respectively, and θ_w represents the weak mixing angle, $\sin^2 \theta_w \simeq 0.23$. The induced electric charge is generally irrational, and it can be relatively large if the hidden $U(1)_H$ is more strongly coupled than the electromagnetic coupling, i.e., $e_H > e_{\text{EM}}$ and if the kinetic mixing is large, $\chi = \mathcal{O}(0.1)$.

To be concrete, let us consider a variant of the viltksmodel, where a complex scalar Φ is coupled to n_q vector-like extra quarks (D, \bar{D}) and n_ψ hidden fermions ($\psi, \bar{\psi}$);

$$-\mathcal{L} = y_q \Phi \sum_{i=1}^{n_q} \bar{D}_i D_i + y_\psi \Phi \sum_{i=1}^{n_\psi} \bar{\psi}_i \psi_i + h.c., \quad (2.8)$$

where the subscript i denotes flavor of the extra quarks and hidden fermions. The charge assignment of these extra matter fields is given in Table 1.^{#4} Here, we focus on the case in which ψ and $\bar{\psi}$ are SM gauge singlets in the original basis and have charges q_H and $-q_H$ under $U(1)_H$, respectively. Then the hypercharge of ψ_i is induced solely by the kinetic mixing as in Eq. (2.7).

We assume that Φ develops a non-zero expectation value in the vacuum,

$$\Phi = \frac{f+s}{\sqrt{2}} e^{i\phi/f}, \quad (2.9)$$

^{#4} One may impose an approximate global $U(1)$ symmetry to ensure the above interaction [6].

	D_i	\bar{D}_i	ψ_i	$\bar{\psi}_i$
SU(3)	3	$\bar{3}$	1	1
U(1) _Y	a	$-a$	$0(q_{\text{eff}})$	$0(-q_{\text{eff}})$
U(1) _H	0	0	q_H	$-q_H$

Table 1: Charge assignment of the extra fermions in the original (canonical) basis.

where s and ϕ denote the radial and phase degrees of freedom, respectively, and f is the decay constant. Then, the extra matter fields have masses of $y_{q,\psi}f/\sqrt{2}$. While the diphoton excess can be explained by either s or ϕ , we will focus on ϕ in the following analysis. Our results can be straightforwardly applied to s except for a possibly large branching fraction of s decaying into a pair of ϕ .^{#5} Hereafter, we assume the mass of ϕ , denoted as m_ϕ , to be 750 GeV.

The field ϕ can decay to gluons and photons via 1-loop diagram and their decay rates are given by

$$\Gamma(\phi \rightarrow gg) = 8 \left(\frac{\alpha_3}{8\pi f} \right)^2 \frac{m_\phi^3}{\pi} \left| \sum_i \frac{1}{4} A_{1/2}(x_i) \right|^2, \quad (2.10)$$

$$\Gamma(\phi \rightarrow \gamma\gamma) = \left(\frac{\alpha_{\text{EM}}}{8\pi f} \right)^2 \frac{m_\phi^3}{\pi} \left| \sum_i \frac{Q_i^2}{2} A_{1/2}(x_i) \right|^2, \quad (2.11)$$

where α_{EM} and α_3 are the electroweak and strong gauge coupling strength, respectively, Q_i is an electric charge of the i -th particle in the loop, and $x_i \equiv 4m_i^2/m_\phi^2$. The form factor $A_{1/2}$ is given by

$$A_{1/2}(x) = 2x \arcsin^2(1/\sqrt{x}) \quad \text{for } x \geq 1, \quad (2.12)$$

which satisfies $A_{1/2}(\infty) = 2$. Then, the production cross section for $pp \rightarrow \phi \rightarrow \gamma\gamma$ is estimated as

$$\begin{aligned} \sigma(pp \rightarrow \phi + X) \text{Br}(\phi \rightarrow \gamma\gamma) &\simeq K \cdot \frac{\pi^2}{8m_\phi} \frac{1}{s} \Gamma(\phi \rightarrow gg) \text{Br}(\phi \rightarrow \gamma\gamma) C_{gg}, \\ C_{gg} &\equiv \int_0^1 dx_1 \int_0^1 dx_2 f_g(x_1) f_g(x_2) \delta(x_1 x_2 - m_\phi^2/s), \end{aligned} \quad (2.13)$$

where K denotes the K-factor, $\sqrt{s} = 13$ TeV and $m_\phi = 750$ GeV. Taking the factorization scale to be $0.5m_\phi$, $C_{gg} \approx 1904$ using MSTW2008NNLO [29]. Since $\Gamma(\phi \rightarrow gg)$ is much larger than $\Gamma(\phi \rightarrow \gamma\gamma)$ in the parameter region of our interest, we use an approximation, $\Gamma(\phi \rightarrow$

^{#5} The decay of s into a pair of ϕ can be suppressed by introducing Φ_1 and Φ_2 with opposite PQ charges, if they respect an approximate Z_2 exchange symmetry, $\Phi_1 \leftrightarrow \Phi_2$.

$gg)\text{Br}(\phi \rightarrow \gamma\gamma) \approx \Gamma(\phi \rightarrow \gamma\gamma)$, resulting in

$$\sigma(pp \rightarrow \phi + X)\text{Br}(\phi \rightarrow \gamma\gamma) \approx K \cdot 7.2 \text{ fb} \left(\frac{\Gamma(\phi \rightarrow \gamma\gamma)}{10^{-3} \text{ GeV}} \right). \quad (2.14)$$

In a simple case of $a = 0$, the cross section is calculated as

$$\sigma(pp \rightarrow \phi + X)\text{Br}(\phi \rightarrow \gamma\gamma) \simeq 3.5 \text{ fb} \left(\frac{f}{800 \text{ GeV}} \right)^{-2} \left(\frac{K}{1.5} \right) \left(\frac{q_{\text{eff}}^2 n_\psi}{4} \right)^2, \quad (2.15)$$

where we assume $A_{1/2}(x_\psi) \simeq 2$ and the K-factor is estimated as $K \approx 1.5$ (see e.g. [5]). In the case of $a \neq 0$, $q_{\text{eff}}^2 n_\psi$ should be replaced with $3a^2 n_q + q_{\text{eff}}^2 n_\psi$ in the above equation.

In our scenario the pseudoscalar ϕ decays into other channels. The ratios of the decay of ϕ into $\gamma\gamma$, γZ , ZZ , $\gamma\gamma'$, $Z\gamma'$, and $\gamma'\gamma'$ are given by

$$\begin{aligned} & \gamma\gamma : \gamma Z : ZZ : \gamma\gamma' : Z\gamma' : \gamma'\gamma' \\ & \simeq 1 : 2 \tan^2 \theta_w : \tan^4 \theta_w : 2 \frac{\alpha_H}{\alpha_{\text{EM}}} \left(\frac{k_{\text{mix}}}{k} \right)^2 : 2 \frac{\alpha_H}{\alpha_{\text{EM}}} \left(\frac{k_{\text{mix}}}{k} \right)^2 \tan^2 \theta_w : \left(\frac{\alpha_H k_H}{\alpha_{\text{EM}} k} \right)^2, \end{aligned} \quad (2.16)$$

with

$$k = 3a^2 n_q + q_{\text{eff}}^2 n_\psi, \quad (2.17)$$

$$k_H = q_H^2 n_\psi, \quad (2.18)$$

$$k_{\text{mix}} = q_{\text{eff}} q_H n_\psi. \quad (2.19)$$

Here we have dropped the phase space factor. As the simplest realization of our scenario, let us focus on the case of $a = 0$. Then, the ratios of the branching fraction of $\phi \rightarrow \gamma\gamma'$ and $\phi \rightarrow Z\gamma'$ to that of $\phi \rightarrow \gamma\gamma$ are

$$\frac{\text{Br}(\phi \rightarrow \gamma\gamma')}{\text{Br}(\phi \rightarrow \gamma\gamma)} = 2 \frac{1 - \chi^2}{\chi^2} \frac{1}{\cos^2 \theta_w} \simeq 2.6 \frac{1 - \chi^2}{\chi^2}, \quad (2.20)$$

$$\frac{\text{Br}(\phi \rightarrow Z\gamma')}{\text{Br}(\phi \rightarrow \gamma\gamma)} = 2 \frac{1 - \chi^2}{\chi^2} \frac{\tan^2 \theta_w}{\cos^2 \theta_w} \simeq 0.78 \frac{1 - \chi^2}{\chi^2}. \quad (2.21)$$

One can relate the production cross section for events with a photon and a hidden photon to that for diphoton events as

$$\sigma(pp \rightarrow \phi \rightarrow \gamma'\gamma) \simeq 10 \text{ fb} \frac{1 - \chi^2}{\chi^2} \left(\frac{\sigma(pp \rightarrow \phi \rightarrow \gamma\gamma)}{4 \text{ fb}} \right). \quad (2.22)$$

The events with a photon and missing transverse momentum have been searched for at the LHC, and their production cross section is constrained to be below 17.8 fb (95%CL) by the ATLAS experiment at $\sqrt{s} = 13 \text{ TeV}$ [13]. This sets a *lower* limit on the kinetic mixing as

$$\chi \gtrsim 0.6, \quad (2.23)$$

where we have set $\sigma(pp \rightarrow \phi \rightarrow \gamma\gamma) = 4 \text{ fb}$. The lower bound on χ is relaxed if $a \neq 0$.

Using the bound on the kinetic mixing, one can derive a constraint on f to explain the diphoton excess,

$$f \gtrsim 800 \text{ GeV} \left(\frac{n_\psi q_H^2 \alpha_H}{0.08} \right) \left(\frac{K}{1.5} \right)^{1/2} \left(\frac{\sigma(pp \rightarrow \phi \rightarrow \gamma\gamma)}{3.5 \text{ fb}} \right)^{1/2}. \quad (2.24)$$

Thus, one can explain the diphoton excess by introducing a single vector-like hidden lepton charged under $U(1)_H$ with a kinetic mixing with $U(1)_Y$, while satisfying the experimental bound on events with a photon and missing transverse momentum. In particular, neither many vector-like matter fields nor large hypercharge in the original basis is needed.

Since fractionally charged particles may be produced via Drell-Yan process, their charges and masses are constrained by the LHC experiment. The ATLAS and CMS collaborations put the upper bound on the mass of stable particle with an electric charge of $(1-2)|e_{\text{EM}}|$ [23, 24]. One can evade the constraint for $m_\psi \gtrsim 600-700 \text{ GeV}$ with $|q_{\text{eff}}| = 1-2$. If \bar{D}_i mixes with a SM quark for $a = 1/3$, it decays into e.g. Higgs and a quark, avoiding constraints from R -hadron searches.^{#6}

3 A model with gauge coupling unification

In this section, we propose a model consistent with the gauge coupling unification. The RG flow of gauge coupling constants is modified by the presence of the large kinetic mixing. In addition, the normalization of $U(1)_Y$ gauge coupling is affected by the large kinetic mixing [see Eq. (2.4)], where we require that the SM gauge couplings are unified in the original basis.

We introduce N_5 pair of $SU(5)$ complete multiplets as

$$- \mathcal{L} = y_D \sum_i^{N_5} \Phi D_i \bar{D}_i + y_L \sum_i^{N_5} \Phi L_i \bar{L}_i + y_\psi \sum_i^{n_\psi} \Phi \psi_i \bar{\psi}_i + h.c., \quad (3.25)$$

where D_i and \bar{L}_i (\bar{D}_i and L_i) consist of $SU(5)$ multiplets transforming **5** ($\bar{\mathbf{5}}$) representation. The $U(1)_Y$ charges of D_i and \bar{L}_i are $-1/3$ and $1/2$, respectively, and they are singlet under the $U(1)_H$ gauge group. Here, ψ_i and $\bar{\psi}_i$ are only charged under $U(1)_H$ as noted in Table 1.

To calculate the RG flow, it is convenient to write the coupling $q_{\text{eff}} g_Y$ as $q_H g_{\text{mix}}$. Thus, in the canonical basis, a matter field Ψ , which collectively denotes D , L and ψ , have an interaction with

$$\mathcal{L} = \bar{\Psi} \gamma^\mu [e_H q_H A_{H\mu} + (g_Y q_Y + g_{\text{mix}} q_H) A_{Y\mu}] \Psi, \quad (3.26)$$

^{#6} If \bar{D}_i mixes with the bottom quark, the lower bound on its mass is severe as $\sim 700 \text{ GeV}$ [25, 26]. However, the bound is much weaker in the case that it mixes with a light quark [27].

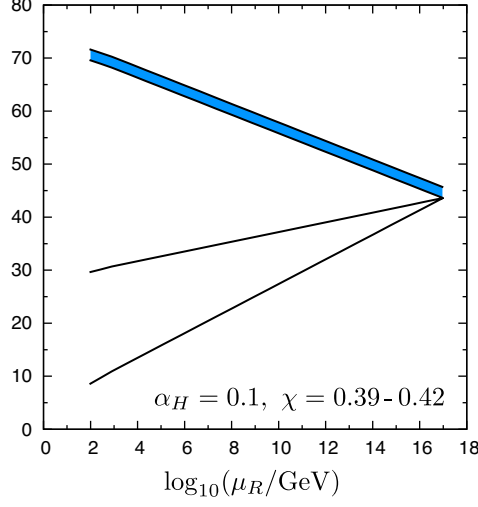


Figure 1: The running of the gauge couplings $(\alpha_1'^{-1}, \alpha_2^{-1}, \alpha_3^{-1})$ in the original basis from top to bottom. Here, $q_H = 1$, $n_\psi = 1$, $N_5 = 1$, $\alpha_3(m_Z) = 0.1185$ and $m_t(\text{pole}) = 173.34$ GeV. The masses of the extra matter fields are taken to be 800 GeV. We take $\alpha_H = 0.1$ and $\chi = 0.39-0.42$ at $\mu_R = m_Z$.

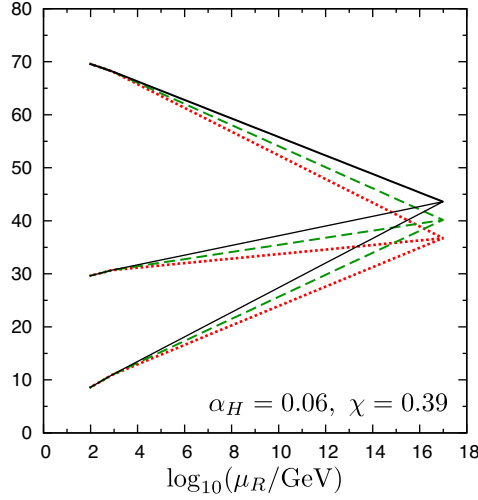


Figure 2: The running of the gauge couplings $(\alpha_1'^{-1}, \alpha_2^{-1}, \alpha_3^{-1})$ in the original basis for $N_5 = 1, 2$ and 3 . The solid, dashed and dotted lines show the RG flow of the gauge couplings for $N_5 = 1, 2$ and 3 , respectively. Here, $\alpha_H = 0.06$ and $\chi = 0.39$ are taken at $\mu_R = m_Z$. The other parameters are the same as in Fig. 1.

where g_{mix} is identified with $-e_H\chi/\sqrt{1-\chi^2}$ at the scale where Eq.(2.1) is defined. Then, the RG equations are [30]

$$\begin{aligned}\frac{dg_Y}{dt} &= \frac{g_Y}{16\pi^2} (b_Y g_Y^2 + b_H g_{\text{mix}}^2), \\ \frac{de_H}{dt} &= \frac{e_H}{16\pi^2} (b_H e_H^2), \\ \frac{dg_{\text{mix}}}{dt} &= \frac{g_{\text{mix}}}{16\pi^2} (b_Y g_Y^2 + b_H g_{\text{mix}}^2 + 2b_H e_H^2),\end{aligned}\tag{3.27}$$

with

$$\begin{aligned}b_H &= \frac{4}{3}n_\psi q_H^2, \\ b_Y &= \frac{41}{6} + \frac{10}{9}N_5.\end{aligned}\tag{3.28}$$

The coefficients of beta functions for $\text{SU}(2)_L$ and $\text{SU}(3)_c$ are given by

$$\begin{aligned}b_2 &= -\frac{19}{6} + \frac{2}{3}N_5, \\ b_3 &= -7 + \frac{2}{3}N_5.\end{aligned}\tag{3.29}$$

Here, $t = \log \mu_R$, where μ_R is a renormalization scale.

We plot the RG flow of couplings α_3 , α_2 , and α'_1 [$\equiv (1-\chi^2)\alpha_1$] in Fig. 1 and 2, where we take the $\text{SU}(5)$ normalization of $\alpha_1 \equiv (5/3)\alpha_Y$ and assume $q_H = 1$ and $n_\psi = 1$. In Fig. 1, we have varied the kinetic mixing slightly, which is represented by the blue band. We find that the SM gauge couplings are unified at the energy scale of order 10^{17} GeV, which is consistent with the null result of proton decay. The relevant RGEs are given by ^{#7}

$$\frac{d\alpha_1'^{-1}}{dt} = -\frac{b_Y}{2\pi} \left(\frac{3}{5}\right), \quad \frac{d\alpha_2^{-1}}{dt} = -\frac{b_2}{2\pi}, \quad \frac{d\alpha_3^{-1}}{dt} = -\frac{b_3}{2\pi}.\tag{3.30}$$

Apparently, the running of α'_1 does not depend on $\alpha_H(m_\psi)$. In Fig. 2, we also show the RG flow of the gauge couplings for $N_5 = 1, 2$ and 3. One can see that the unification point at around 10^{17} GeV is independent of N_5 . The hidden gauge coupling e_H remains perturbative up to the GUT scale:

$$\alpha_H(10^{17} \text{ GeV}) = \alpha_H(m_\Psi) \left[1 - \frac{\alpha_H(m_\Psi)}{2\pi} \ln \frac{10^{17} \text{ GeV}}{m_\Psi} \right]^{-1} \approx 0.21 (0.09),\tag{3.31}$$

for $\alpha_H(m_\Psi) = 0.10 (0.06)$.

^{#7} Even if there exist fields which have both $\text{U}(1)_Y$ and $\text{U}(1)_H$ charges, the form of the RGE, $\frac{d\alpha_1'^{-1}}{dt} = -\frac{b_Y}{2\pi} \left(\frac{3}{5}\right)$, does not change.

With the extra SU(5) multiplets, the cross section of the diphoton signal is modified such as $k \rightarrow k + N_5$, $n_q = N_5$, and $a = -1/3$ due to the presence of the GUT multiplet. Taking $q_H^2 n_\psi = 1$, $N_5 = 1$, $\chi = 0.40$ and $\alpha_H = 0.1$, the cross section is estimated as

$$\sigma(pp \rightarrow \phi + X) \text{Br}(\phi \rightarrow \gamma\gamma) \approx 3.4 \text{ fb} \left(\frac{f}{800 \text{ GeV}} \right)^{-2}. \quad (3.32)$$

Here, the masses for the extra matter fields are taken as $m_\psi = m_{D'} = 650 \text{ GeV}$ and $m'_L = 450 \text{ GeV}$, where $m_{D'} = y_D f / \sqrt{2}$ and $m_{L'} = y_L f / \sqrt{2}$. In the case $N_5 = 2$, the same cross section of Eq. (3.32) is obtained with the smaller α_H and larger f : $\alpha_H = 0.06$ and $f = 980 \text{ GeV}$ with the other parameters being the same as the previous case.

The ratio of the branching fraction of $\phi \rightarrow \gamma\gamma'$ to that of $\phi \rightarrow \gamma\gamma$ is suppressed due to the contribution to the latter process from the GUT multiplet:

$$\frac{\sigma(pp \rightarrow \phi \rightarrow \gamma\gamma')}{\sigma(pp \rightarrow \phi \rightarrow \gamma\gamma)} \approx \frac{n_\psi^2 q_{\text{eff}}^4}{(N_5 4/3 + n_\psi q_{\text{eff}}^2)^2} \frac{1 - \chi^2}{\chi^2} \frac{2}{\cos^2 \theta_w}. \quad (3.33)$$

As a result, we find that the present constraint on the events with a photon and missing transverse momentum can be suppressed compared to the case without extra SU(5) multiplets. For a set of the parameters consistent with the GUT, $N_5 = 1$, $n_\psi q_{\text{eff}}^2 = 1.92$, $\chi = 0.40$ ^{#8} and $\sigma(pp \rightarrow \phi \rightarrow \gamma\gamma) = 3.4 \text{ fb}$, we have $\sigma(pp \rightarrow \phi \rightarrow \gamma\gamma') \approx 16.1 \text{ fb}$. Obviously, larger N_5 leads to a weaker constraint: for $N_5 = 2$, $n_\psi q_{\text{eff}}^2 = 1.15$ and $\chi = 0.40$, the relevant cross section is as small as $\sigma(pp \rightarrow \phi \rightarrow \gamma\gamma') \approx 4.2 \text{ fb}$.

Finally, let us comment on a possible generation of the large mixing. In $\text{SU}(5)_{\text{GUT}} \times \text{U}(1)_H$ model, we may have the following operator:

$$\frac{1}{M_*} \text{Tr}(\Sigma_{24} F_{\mu\nu}^5) F_H^{\mu\nu}, \quad (3.34)$$

where $F_{\mu\nu}^5$ is a gauge field strength of $\text{SU}(5)_{\text{GUT}}$; M_* is a cut-off scale, and Σ_{24} is a GUT breaking Higgs with $\langle \Sigma_{24} \rangle = \text{diag}(2, 2, 2, -3, -3) v_{\text{GUT}}$. (Here, $v_{\text{GUT}} \sim 10^{17} \text{ GeV}$). Therefore, if M_* is somewhat close to v_{GUT} , the large mixing between $\text{U}(1)_Y$ and $\text{U}(1)_H$ can be generated via the above high dimensional operator. That said, it is fair to admit that obtaining such a large kinetic mixing is highly nontrivial in a context of gauge coupling unification. This is because we can similarly write down the following operator:

$$\frac{1}{M_*} \text{Tr}(\Sigma_{24} F_{\mu\nu}^5 F^{5\mu\nu}), \quad (3.35)$$

which could generate a large threshold correction, preventing the gauge couplings from precise unification. We note however that the relative size of these operators Eqs. (3.34) and (3.35) depends on details of UV physics. ^{#9}

^{#8} At the GUT scale 10^{17} GeV , $\chi \sim 0.6 - 0.8$, depending on e_H .

^{#9} For instance, we can consider interactions:

$$\mathcal{L} = \lambda_5 f_{\mathbf{5}}^l \Sigma_{24} f_{\mathbf{5}}^l + M_5 f_{\mathbf{5}}^l f_{\mathbf{5}}^l + h.c., \quad (3.36)$$

4 Cosmology

In this section, we explain cosmology of our model. The hidden photon decouples from the SM sector at a temperature around $m_\psi/10$ and contributes to the energy density of the Universe as dark radiation, whose amount can be measured by future observations of CMB temperature fluctuations. Since the fields ψ_i and $\bar{\psi}_i$ have fractional charges of $U(1)_Y$, the lightest ones are absolutely stable. On the other hand, extra quarks (and leptons) mix with SM quarks (leptons), and are not stable as mentioned earlier. The abundance of the fractionally charged particles, ψ_i and $\bar{\psi}_i$, are severely constrained by various experiments and observations. We provide some possibilities to evade these constraints in Sec. 4.4.

4.1 Dark radiation

Since the hidden $U(1)_H$ is not broken in our model, we predict hidden photon as well as the fractionally charged particles. When the temperature is higher than the mass of ψ , the $U(1)_H$ gauge boson as well as ψ are in thermal equilibrium with the SM plasma. Even after the temperature decreases to m_ψ , the $U(1)_H$ gauge boson may interact with visible photon via photon-photon scatterings. Here we quote the low-energy scattering cross section of visible photons:

$$\frac{d\sigma(\gamma\gamma \rightarrow \gamma\gamma)}{d\Omega} = \frac{139}{(180\pi)^2} \alpha_{\text{EM}}^4 \frac{\omega^6}{m_e^8} (3 + \cos^2\theta)^2, \quad (4.1)$$

where m_e is the electron mass and ω is the energy of each colliding photon in the frame in which the total momentum vanishes and θ is a scattering angle. We expect that scatterings between visible photon and hidden photon is roughly given by Eq. (4.1) with the replacement of $\alpha_{\text{EM}}^4 \rightarrow q_{\text{eff}}^2 \alpha_{\text{EM}}^2 \alpha_H^2$ and $m_e \rightarrow m_\psi$ with an additional $\mathcal{O}(1)$ coefficient. We find that this kind of interaction decouples at a temperature of order $m_\psi/10$ for $m_\psi = \mathcal{O}(1)$ TeV. However, the fractionally charged particles ψ may be still in thermal equilibrium with both the SM plasma and the hidden sector via Compton scatterings. The Compton scattering between ψ and $U(1)_H$ gauge boson is decoupled at a temperature satisfying $\sigma'_T n_\psi/H \sim 1$, where σ'_T

where $f_{\mathbf{5}}^l$ has a $U(1)_H$ charge of $q_{H,f}$, and $l = 1 \dots N_f$. Then, after integrating out $f_{\mathbf{5}}^l$ and $f_{\mathbf{5}}^l$, the mixing term is generated as

$$\mathcal{L} \sim \frac{\lambda_5 g_5 e_H q_{H,f} N_f}{16\pi^2 M_f} \text{Tr}(\Sigma_{24} F_{\mu\nu}^5) F_H^{\mu\nu}, \quad (3.37)$$

where $M_f \sim \lambda_5 v_{\text{GUT}} + M_5$. For $\lambda_5 \sim 4\pi$, $M_f \approx v_{\text{GUT}}$, and $(q_{H,f} N_f) \sim 10$, the mixing becomes $\mathcal{O}(1)$. In addition, Eq. (3.35) is generated as

$$\mathcal{L} \sim \frac{\lambda_5 g_5^2 N_f}{16\pi^2 M_f} \text{Tr}(\Sigma_{24} F_{\mu\nu}^5 F_5^{\mu\nu}). \quad (3.38)$$

The relative size of these operators depends on $q_{H,f} e_H$, so that we can suppress the coupling constant of the latter operators with fixing that of the former operator. We thank an anonymous referee for pointing out this issue.

is the Thomson scattering rate given by $8\pi\alpha_H^2/3m_\psi^2$. When the number density of n_ψ is determined by the thermal relic density [see Eq. (4.5)], the combination is rewritten as

$$\frac{\sigma'_T n_\psi(T)}{H(T)} \simeq \frac{8}{3} \frac{T_f}{T}, \quad (4.2)$$

where T_f is the freezeout temperature of ψ ($\approx m_\psi/25$). Thus the $U(1)_H$ decouples from the SM plasma at a temperature of order $m_\psi/70$.

After the $U(1)_H$ gauge boson decouples from the SM thermal plasma, its energy density contributes to the expansion of the Universe as dark radiation. Its amount is conventionally expressed by the effective neutrino number ΔN_{eff} , which is calculated as

$$\Delta N_{\text{eff}} = \frac{8}{7} \left(\frac{g_*(T_D)}{43/4} \right)^{-4/3}, \quad (4.3)$$

where $g_*(T_D)$ is the effective relativistic degrees of freedom at the decoupling temperature T_D (see, e.g., Refs. [31–33]). It is given as $g_*(T_D) \simeq 103.9$ for $T_D = 200\text{GeV}$, $g_*(T_D) \simeq 103.5$ for $T_D = 100\text{GeV}$, $g_*(T_D) \simeq 97.4$ for $T_D = 50\text{GeV}$, and $g_*(T_D) \simeq 86.2$ for $T_D = 10\text{GeV}$, which imply that the effective neutrino number is about $0.054 - 0.071$. The Planck data combined with the observation of BAO puts the constraint $N_{\text{eff}} = 3.15 \pm 0.23$ [34], which is consistent with the value predicted in the standard model ($N_{\text{eff}} = 3.046$) and our prediction. The deviation from the standard value will be observed by the ground-based Stage-IV CMB polarization experiment CMB-S4, which measures N_{eff} with a precision of $\Delta N_{\text{eff}} = 0.0156$ within 1σ level [35] (see also Ref. [36]).

4.2 Primordial abundance of charged particles

When the reheating temperature of the Universe is higher than the freezeout temperature of ψ , its thermal relic abundance is determined as

$$\Omega_\psi h^2 \approx \frac{5.0 \times 10^{-27} \text{ cm}^3 \text{ s}^{-1}}{\langle \sigma v \rangle}, \quad (4.4)$$

where the annihilation cross section is given by

$$\langle \sigma_\psi v \rangle \simeq \frac{\pi \alpha_H^2}{m_\psi^2} + N \frac{q_{\text{eff}}^4 \pi \alpha_Y^2}{m_\psi^2}. \quad (4.5)$$

The second term includes the annihilation into the SM particles and the prefactor N is given by $N \simeq 1 + (5 + 1/8)/q_{\text{eff}}^2$ for $m_\psi \gg \mathcal{O}(100\text{GeV})$. Below we neglect the annihilation process via the EW force because its coupling constant is much smaller than that of hidden $U(1)_H$ ($\alpha_H \approx 0.1$). Note that the annihilation cross section increases by a factor of 1.2 when we assume $q_{\text{eff}} = 2$ and take into account the annihilation into the SM particles. For typical parameters, their abundance is given by

$$\Omega_\psi h^2 \approx 0.013 \left(\frac{m_\psi}{1 \text{ TeV}} \right)^2 \left(\frac{\alpha_H}{0.1} \right)^{-2}. \quad (4.6)$$

Fractionally charged particles may affect CMB temperature fluctuations, so that precise measurement of CMB temperature fluctuations provides an upper bound on their abundance. However, all of the previous works focused on the case of millicharged particles. Therefore their result cannot directly apply to our case, where the electric charge of ψ is of order unity. Still, we expect that the abundance of these particles should be less than of order the uncertainty of baryon abundance determined by the Planck experiment. As discussed in Refs. [20, 21], their constraint comes from the fact that the exotic charged particles are tightly coupled with the plasma before the recombination epoch but the Compton scattering process is neglected due to its small electric charge. In particular, the constraint given in Ref. [21] is based mainly on the fact that increasing the number density of millicharged particles results in decreasing that of baryons, which results in decreasing that of electrons by the neutrality condition of the Universe. As a result, the Silk damping scale becomes larger compared with the case without millicharged particles. In our case, the tight coupling condition is satisfied due to the large electric charge while the Compton scattering process is neglected due to the suppression of the cross section by the large mass of ψ . Therefore we can apply their result to our case with $\mathcal{O}(1)$ electric charge. Thus we require [21]

$$\Omega_\psi h^2 \lesssim 0.001. \quad (4.7)$$

Another constraint comes from the observation of Li abundance, which is marginally consistent with the prediction of the BBN theory without fractionally charged particles. When electrically charged particles are abundant in the BBN epoch, they form a bound state with ^4He , which leads to an efficient production of ^6Li via a photon-less thermal production process [17, 18]. The enhancement of ^6Li production originates mainly from the fact that the Bohr radius of the bound state is much shorter than the wavelength of emitted photon in the standard BBN theory. Since the Bohr radius is determined by the mass of nucleus and the charges of bounded particles, their results do not change by many orders of magnitude even in the case with fractionally charged particle with $\mathcal{O}(1)$ electric charge. Thus we quote their results [19]:

$$\frac{n_\psi}{n_b} \lesssim 10^{-5-6}, \quad (4.8)$$

where n_b is the baryon number density. This constraint is severer than the one coming from the observations of CMB temperature fluctuations. We discuss how to evade these constraints in Sec. 4.4.

4.3 Present abundance of charged particles

Next, we consider an era after the solar system and the Earth form, following Ref. [37]. The number density of fractionally charged particles in bulk matter (i.e., in the Earth or solar system) is different from that given in Eq. (4.6) because of their electrical interaction with matter, which results in efficient annihilation in bulk matter. However, the annihilation is

not so efficient that we cannot avoid severe constraints by the null results of searches of fractionally charged particles in bulk matter.

Fractionally charged particles in the Earth are more dense than their average density in the Universe because they behave like baryons due to their electric charge. In addition, the annihilation cross section is enhanced by the Sommerfeld enhancement effect in a low terrestrial temperature:

$$\langle \sigma_\psi v \rangle_{\text{SF}} = S(\eta) \langle \sigma_\psi v \rangle, \quad (4.9)$$

where $S(\eta)$ is a Sommerfeld enhancement factor given by

$$S(\eta) \equiv \frac{\eta}{1 - e^{-\eta}}. \quad (4.10)$$

The parameter η is defined by

$$\eta \equiv 2\pi \frac{\alpha_H}{\beta}, \quad (4.11)$$

where β is the velocity of the fractionally charged particle in a low terrestrial temperature of order 300 K. Since the time scale is the age of the Earth, which is of order 4.5 Gyr ($\equiv t_E$), the annihilation reduces the abundance of fractionally charged particles to the amount of

$$\left(\frac{n_\psi}{n_B} \right) \simeq \frac{1}{n_B \langle \sigma_\psi v \rangle_{\text{SF}} t_E}. \quad (4.12)$$

Using the number density of baryons in bulk matter of $n_B \simeq 6.4 \times 10^{23} \text{ cm}^{-3}$, we obtain

$$\left(\frac{n_\psi}{n_B} \right) \simeq 8.5 \times 10^{-24} \left(\frac{m_\psi}{1 \text{ TeV}} \right)^{3/2} \left(\frac{\alpha_H}{0.1} \right)^{-3}. \quad (4.13)$$

One may wonder that negatively fractionally charged particles capture protons and/or Heliums and form positive exotic ions, which cannot annihilate with anti-particles due to the electrical repulsion of Coulomb force [37]. However, in our model, the annihilation occurs due to the hidden $U(1)_H$ gauge interaction, which is much stronger than the electric force, so that annihilation cannot be prevented by the visible Coulomb force. Therefore the abundance of fractionally charged particles in the Earth is given by Eq. (4.13).

Fractionally charged particles can be observed by searching in bulk matter if they are trapped in rigid matter or water. Most of the searches of fractionally charged particles put constraints on the abundance of particles with a fractional charge in the interval of $[n + 0.2, n + 0.8]$ where n is any integer [38, 39] (see Ref. [40] for a review). In the recent paper of Ref. [41], however, they provided a constraint which is less stringent but is applicable to broader range of charges by using optically levitated microspheres in high vacuum. They also claimed that the previous works can constrain the abundance of particles with smaller charges by assuming the abundance of negative fractionally charged particles. This is because there can be multiple fractionally charged particles in each sample when their number

density is sufficiently large. As a result, the total charge in each sample is the summation of charges of those particles, which can be larger than about 0.2 and can be detectable by those experiments. Their results indicate that the abundance of fractionally charged particles has to be 15-23 order of magnitude less than that of baryons, depending on their charge. For example, for a fractional charge in the interval of $[n + 0.2, n + 0.8]$ abundance per nucleon should be less than 10^{-23} , and for a fractional charge in the interval of $[n + 0.1, n + 0.2]$ and $[n + 0.8, n + 0.9]$ abundance per nucleon should be less than 10^{-20-21} . The result of Eq. (4.13) is consistent with this upper bound. However, we should also consider the abundance of fractionally charged particles in the interstellar medium (ISM). Its number density in the ISM is much less than that in the Earth, so that the annihilation is inefficient as discussed in Ref. [37]. This may imply that the searches of fractionally charged particles in meteorites (e.g., the work of Ref. [39]) exclude our model though the evolutionary history of meteoritic material is uncertain. In addition, the calculation of Eq. (4.13) does not take into account the flux of fractionally charged particles from the outer region of the Earth. Since they are abundant in the outer region, fractionally charged particles may fall into the Earth just like cosmic rays. Therefore the constraints coming from the search in bulk matter may exclude the scenario that fractionally charged particles survive at present. In the next subsection, we provide some mechanisms to evade those constraints.

4.4 Possibilities to eliminate unwanted relics

We can consider a scenario in which the unwanted charged particles are never produced after inflation, which requires that the maximal temperature of the Universe after inflation is much lower than the mass of ψ .^{#10} Alternatively, one may consider huge late-time entropy production by thermal inflation. In this case, the hidden photon is also diluted, so that the dark radiation is absent [see Eq. (4.3)].

Another way to evade the constraints is to enhance their annihilation rate by a strong interaction. We may introduce an additional $U(1)_{H2}$ gauge symmetry under which the fields ψ_i are charged [43, 44]. We also introduce a scalar monopole that develops an expectation value of order 1 TeV to break $U(1)_{H2}$ spontaneously.^{#11} As a result, the fields ψ_i are connected by the physical string due to the dual Meisner effect [45], so that they soon annihilate with each other after the spontaneous symmetry breaking (SSB). Therefore the fields ψ_i are absent after the SSB, so that we can evade the constraint coming from the searches of bulk matter as well as that coming from the observations of CMB temperature fluctuations.^{#12} We assume that there is no additional mixing among $U(1)_Y$ and $U(1)_{H2}$ so

^{#10} In Ref. [42], one of the authors (M.Y.) investigated the thermalization process of inflaton decay products and found that the maximal temperature of the Universe after inflation can be much lower than the one expected in the literature due to the delay of thermalization. It was found that the maximal temperature can be less than 100 GeV.

^{#11} Since ψ_i has both charges of $U(1)_H$ and $U(1)_{H2}$, an operator $g^{\rho\sigma}|\phi_M|^2 F_{\mu\nu}^H F_{\rho\sigma}^H$ arises where ϕ_M is a scalar monopole. Therefore, it may be possible to test a TeV scale photon-photon collider depending on the sizes of the operator and χ .

^{#12} The $U(1)_{H2}$ gauge theory may be conformal in the presence of monopole as well as electrons. Thus we

that our calculations given in the previous sections are not changed.^{#13} In this case, when ψ and $\bar{\psi}$ are produced via the Drell-Yan process at the collider experiment, they form a bound state decaying into $\gamma\gamma$, $\gamma\gamma'$ and $\gamma'\gamma'$. Also, when the center of energy available is sufficiently large, mesonization occurs and four-photon signals will be observed due to the subsequent annihilation of the mesons ($\psi\bar{\psi}$).

When the number of flavor of the hidden particle is larger than unity, we can predict a long-lived neutral particle that can be a candidate for DM. We consider ψ_i with $i = 1, 2$ and assume that the flavour is not mixed so that the $\bar{\psi}_1\psi_2$ bound state (which we denote π_{DM}) does not annihilate after the SSB of $U(1)_{H2}$. Since this bound state is neutral and stable, it can be DM. When their masses are larger than v , the relic abundance of the fields ψ_i is determined by their annihilation rate [see Eq. (4.6)]. Then they are attached by the physical string after the SSB. The other bound states (e.g., $\bar{\psi}_1\psi_1$) annihilate into visible photons. As a result, the relic density of ψ is given by

$$n_\psi \sim \frac{n_{\psi_1}n_{\psi_2}}{n_{\psi_1} + n_{\psi_2}}. \quad (4.14)$$

This is consistent with the observed DM abundance when the masses of ψ_i are of order 1 TeV.^{#14} The DM can decay when we write a higher dimensional operator of $\psi_1\bar{\psi}_1\psi_1\bar{\psi}_2/M_{\text{pl}}^2$. However, its lifetime is much larger than the present age of the Universe for $\Lambda \lesssim 100$ GeV, so that we expect no astrophysical signal from DM decay.

5 Discussion and conclusion

We have shown that the 750 GeV diphoton excess can be explained by introducing a pair of vector-like quarks and hidden fermions charged under a hidden $U(1)$ gauge symmetry which has a significantly mixing with $U(1)_Y$. The hidden $U(1)$ has a relatively large coupling constant, which may naturally arise from string theory compactifications [47]. Due to the large coupling and kinetic mixing, hidden fermion loops induce a sizable branching fraction of a 750 GeV scalar to diphotons. Notably, the standard model gauge couplings unify around 10^{17} GeV with effects of the kinetic mixing, suggesting the grand unified theory without too rapid proton decay. Obviously, instead of introducing the hidden fermions, one can consider vector-like quarks (and leptons) charged under both the SM gauge symmetry and the hidden $U(1)$ gauge symmetry, which leads to similar results.

need to take care of anomalous dimension of $U(1)_{H2}$ gauge field strength, which may result in the absence of interactions between $U(1)_{H2}$ and $U(1)_H$ [46]. In this paper we neglect this issue by assuming that the conformal coupling constant is not large and the anomalous dimension is sufficiently small.

^{#13} A kinetic mixing between $U(1)_H$ and $U(1)_{H2}$ can be removed by the shift of $U(1)_{H2}$ charge of ψ_i .

^{#14} On the other hand, when their masses are smaller than v , bound states form at the SSB and then their abundance is determined by the subsequent annihilation. The annihilation rate of π_{DM} is estimated as Λ^{-2} where Λ ($\approx 4\pi v$) is the dynamical scale of the confinement. Thus we can account for the observed DM density when the dynamical scale is around the unitarity bound of order 100 TeV. In this case, however, the cross section of the diphoton signal is suppressed and we cannot explain the excess reported by ATLAS and CMS.

Our scenario can be checked by looking for events with a photon and missing transverse momentum. Its cross section is related to that for the diphoton final state through the kinetic mixing, and predicted to be large. Therefore, our scenario can be tested in the near future at the LHC experiment.

So far we have focused on the case in which the phase component, ϕ , is responsible for the diphoton excess. As mentioned earlier, our analysis can be similarly applied to the radial component, s , in Eq. (2.9); in general, s with a mass of 750 GeV decays mainly into a pair of ϕ , each of which decays into $\gamma\gamma$, $\gamma\gamma'$, and $\gamma'\gamma'$, in addition to the loop-induced decays. The photons ($\gamma\gamma$) produced from the decay of ϕ are collimated if the axion is sufficiently boosted, and the two collimated photon pairs may be identified with the diphoton signal in the detector analysis. (See Refs. [7, 48–53] for the collimated photons in association with the diphoton excess.) In our model, the ratio of the branching fraction into the combination of γ and γ' depends on the kinetic mixing. As a result, we may have triphoton and tetraphoton signals, depending on the probability that the collimated photons are identified with a single photon at the detector. If those signals are confirmed by the upcoming LHC data, it would give a smoking gun signature of the diaxions decaying into photons and hidden photons. Further analysis of the above processes is warranted.

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